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Citation: Physics of Plasmas **22**, 091505 (2015); doi: 10.1063/1.4928941
View online: http://dx.doi.org/10.1063/1.4928941
View Table of Contents: http://scitation.aip.org/content/aip/journal/pop/22/9?ver=pdfcov
Published by the AIP Publishing

Cosmic ray confinement and transport models for probing their putative sources

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(Received 20 January 2015; accepted 3 March 2015; published online 27 August 2015)

Recent efforts in cosmic ray (CR) confinement and transport theory are discussed. Three problems are addressed as being crucial for understanding the present day observations and their possible telltale signs of the CR origin. The first problem concerns CR behavior right after their release from a source, such as a supernova remnant. At this phase, the CRs are confined near the source by self-emitted Alfven waves. The second is the problem of diffusive propagation of CRs through the turbulent interstellar medium. This is a seemingly straightforward and long-resolved problem, but it remains controversial and reveals paradoxes. A resolution based on the Chapman-Enskog asymptotic CR transport analysis, that also includes magnetic focusing, is suggested. The third problem is about a puzzling sharp ($\sim 10^{\circ}$) anisotropies in the CR arrival directions that might bear on important clues of their transport between the source and observer. The overarching goal is to improve our understanding of all aspects of the CR's source escape and ensuing propagation through the galaxy to the level at which their sources can be identified observationally. © *2015 AIP Publishing LLC*. [http://dx.doi.org/10.1063/1.4928941]

I. INTRODUCTION

Cosmic rays (CR) have been discovered more than a century ago but the problem of their origin is still with us today. The fundamental obstacle to identification of their possible sources, such as the supernova remnant (SNR) shocks, is the CR "black-box" propagation through the chaotic magnetic field of the galaxy. With a possible exception for the highest energy CR ($\gtrsim 10^{19}$ eV, whose origin is almost certainly extragalactic[1,2]), most of the CRs arrive from random directions saying nothing about the locale of their sources. The more surprising is a sharp ($\sim 10^{\circ}$) CR anisotropy discovered by Milagro[3] with interesting ramifications due to IceCube, ARGO-YBJ,[4] HAWC,[5] and some other instruments. Had it been created in the source, it would have been completely erased en route to the Earth. Therefore, this CR feature is likely to be an imprint of their interaction with the ISM (interstellar medium) or its local environment (heliosphere and surroundings). We will discuss this later. More natural is to start with the CR transport right at their birth place.

SNRs are widely regarded as the most probable source of the bulk of the CRs.[2] The *joint analysis* of the broad band observations of the SNRs and CR background spectra at the Earth should provide the ultimate evidence for this hypothesis. The analysis faces multiple problems. First, the accelerated CRs manifest themselves in SNRs only in form of secondary emission, which is usually difficult to interpret. For example, the super TeV-photons, carefully counted by the atmospheric Cerenkov telescopes, to testify for the accelerated protons colliding with the ambient gas,[6,7] can easily be confused with the inverse Compton (IC) photons upscattered by accelerated electrons. If this is the case, not much weight can be added to the argument for the CRs origin in the SNRs, as electrons comprise only a small fraction ($\sim 1\% - 2\%$) of the CR spectrum. Similarly, in the GeV energy band the emission may come from electron Bremsstrahlung. The key in both cases, however, is a dense gas in the SNR surroundings, often present in form of adjacent molecular clouds (MC). They provide a target for the *pp* reactions with accelerated protons. Photons, produced in this reaction should thus come from MCs, illuminated by the CR protons that have, in turn, escaped from the source. By contrast, the IC electron emission should come from the entire volume they fill, as the low-energy photons (such as CMB) are present everywhere. To use this simple but powerful diagnostic tool for identification of the source of emission detected by modern ground-based instruments and space observatories (such as the Cerenkov telescopes HESS, VERITAS, MAGIC and Fermi-LAT, PAMELA, Agile spacecraft observatories[8–14]), we must have an accurate understanding of the CR propagation from their sources to the adjacent MCs.

The next problem is the subsequent interaction of the CRs with the MC, as their confinement inside the cloud is generally deteriorated due to the collisional damping of the Alfven waves, which otherwise would prevent CRs from spreading further rapidly. In addition, this interaction reveals important clues as to how the spectrum of CRs, illuminating the MC is different from that in the source, most importantly, in form of spectral breaks. This aspect of the CR interaction with MC and their visibility in the gamma-ray band has been discussed in one of the earlier APS Plasma Physics meetings.[15] Here, we will focus on the ensuing propagation of the CR to the Earth and their spectral features that they can acquire both in the source and on the way to us.

Measurements of the CR background spectrum have also advanced significantly. Much progress has been made in isolating different elements in it. One of the most striking results was the ≈ 0.1 difference between the rigidity (momentum to charge ratio) spectral indices of protons and He^{2+}

ions. Such deviations have been apparent for some time, e.g., Ref. 16, but the Pamela spacecraft observatory measured it with a three-digit accuracy in the 100 GeV energy band,[12] which posed a strong challenge to the CR acceleration and propagation models. Indeed, the ultrarelativistic parts of the rigidity spectra must be identical, if protons and He^{2+} ions are accelerated and transported via electromagnetic interactions under identical conditions. Since He^{2+} has a 0.1-flatter spectrum, the difference may be due to its spallation, biased for lower energies.[17] However, such scenario would probably require stretching the model parameters too much.[18] Other interesting possibilities discussed in the literature include the contribution from multiple SNRs of different types with somewhat different CR spectra[12,19–21] and variable p/He^{2+} mix along the shock path.[22,23] A "plasma physics" solution that targets the nonrelativistic phase of acceleration of both species, where the argument of equal rigidity spectra is irrelevant, was suggested in Ref. 24. This explanation is advantageous according to *Occam's razor*, as it relies on the collisionless shock *intrinsic* properties and does not require any of the above special conditions.

Apart from the elemental composition, other spectral signatures, such as the spectral hardening above $E \sim 200$ GeV, have been studied[12,17,25] and provided important clues for the energy dependent CR transport. Stochasticity of CR sources and inhomogeneity of transport through the galaxy is now also included in the models.[26] These are important for understanding the large scale CR anisotropy. The most puzzling aspect of the anisotropy in the CR arrival directions is, in my view, the sharp anisotropy or the so-called Milagro "hot spots" which I address later in this brief review.

The remainder of the paper is organized as follows. In Sec. II, the confinement of CRs released from the source is addressed. In Sec. III, an equation describing diffusive propagation including a hyperdiffusive term is presented and its relation to the so-called "telegrapher" term in the CR transport equation is clarified. In Sec. IV, possible mechanisms for building a sharp CR anisotropy during their propagation from the source are discussed.

II. SELF-CONFINEMENT OF CRs AROUND SNRs

It is widely believed that CRs are accelerated in SNR shocks by the diffusive mechanism (DSA). The backbone of the DSA is a self-confinement of accelerated particles supported by their scattering off magnetic irregularities that particles drive by themselves while streaming ahead of the shock. Logically, this process should also control the ensuing propagation (escape) of CRs, at least until their density drops below the wave instability threshold. At the same time, no consensus has been reached so far as to how CRs escape the accelerator. The dividing lines seem to run across the following issues: (i) does the escape occur isotropically or along the local magnetic field? (ii) does the scattering by the background MHD turbulence control the CR propagation alone or self-excited waves need to be included? (iii) are CRs, that escape SNR, peaked at the highest energy or lower energy CRs escape as well?[22,27–32]

Adhering to the self-confinement idea, we consider the model that explicitly includes the self-excited waves. Moreover, in the regions where magnetic perturbations are weak, i.e., $\delta B^2/B^2 \ll 1$, a field aligned CR transport dominates, as the perpendicular diffusion is suppressed, $\kappa_\perp \simeq (\delta B/B)^2 \kappa_B \ll \kappa_\parallel \simeq \kappa_B (\delta B/B)^{-2}$. Here, κ_B is the Bohm diffusion coefficient $\kappa_B = c r_g/3$ with r_g being the particle gyroradius. Taking into account the condition $(\delta B^2/B^2)_{\rm ISM} \ll 1$, such regime is inevitable outside the source where $\delta B/B \lesssim 1$, as well as at later times of CR propagation, when they are spread over a large volume and the waves are driven weakly. Moreover, the self-confinement of CRs propagating away from the accelerator, as described below, is the continuation of physically the same process long believed to be at work inside the accelerator, as first suggested by Bell.[33] From a mathematical standpoint, our treatment below generalizes Bell's steady state solution, obtained in the shock frame, to the time dependent solution for the CR cloud expanding further out. This being said, we use the following equations that describe the CR diffusion and wave generation self-consistently:[34]

$$\frac{\partial}{\partial t} P_{\rm CR}(p) = \frac{\partial}{\partial z} \frac{\kappa_{\rm B}}{I} \frac{\partial P_{\rm CR}}{\partial z}, \tag{1}$$

$$\frac{\partial}{\partial t} I = -C_{\rm A} \frac{\partial P_{\rm CR}}{\partial z} - \Gamma I, \tag{2}$$

where C_A is the Alfvén velocity. The dimensionless CR partial pressure $P_{\rm CR}$ is used instead of their distribution function $f(p,t)$

$$P_{\rm CR} = \frac{4\pi}{3} \frac{2}{\rho C_{\rm A}^2} v p^4 f, \tag{3}$$

where v and p are the CR speed and momentum, and ρ—the plasma density. The total CR pressure is normalized to $d\ln p$, similar to the wave energy density I

$$\frac{\langle \delta B^2 \rangle}{8\pi} = \frac{B_0^2}{8\pi} \int I(k) d\ln k = \frac{B_0^2}{8\pi} \int I(p) d\ln p.$$

The last relation implies a simplified wave-particle resonance condition, $k r_g(p) = const \sim 1$. Most of the works on CR self-confinement (see Ref. 35 for a review) use equations largely similar to Eqs. (1) and (2), but different assumptions are made regarding geometry of particle escape from the source, the character and strength of the wave damping Γ, and the role of quasilinear wave saturation. A reasonable choice of the damping mechanism is the Goldreich-Shridhar (GS) MHD cascade,[36] which seems to be appropriate in $I \lesssim 1$ regime.[37–39] The damping rate in this case is $\Gamma = C_A \sqrt{k/l}$, where l is the outer scale of turbulence, which may be as large as 100 pc (see, however, Sec. IV). As Γ does not depend on I and can be considered also as coordinate independent, it allows the following ("quasilinear") integral of the system given by Eqs. (1) and (2):

$$P_{\rm CR}(z,t) = P_{\rm CR0}(z) - \frac{\kappa_{\rm B}}{C_{\rm A}} \frac{\partial}{\partial z} \ln \frac{I(z,t)}{I_0(z)}. \tag{4}$$

Here, $P_{CR0}(z)$ and $I_0(z)$ are the initial distributions of the CR partial pressure and the wave energy density (see Ref. 34 for more general treatment). Substituting P_{CR} in Eq. (2), we arrive at the following diffusion equation for I

$$\frac{\partial I}{\partial t}=\frac{\partial}{\partial z}\frac{\kappa_B}{I}\frac{\partial I}{\partial z}-\Gamma I-C_A\frac{\partial P_{CR0}}{\partial z}.$$

Outside of the region where $P_{CR0}\neq 0$, the last term on the r.h.s. is absent, while the second term may be eliminated by replacing $I\exp(\Gamma t)\to I$, $\int_0^t \exp(\Gamma t)dt\to t$. However, if Γ is taken in a GS-form, it is fairly small due to the factor $\sqrt{r_g/l}\ll 1$. We may simply neglect it. The solution for I and $P_{CR}(z,t)$ may be found in an implicit form (see Ref. 34 for details). However, there exists a very accurate convenient interpolation formula that can be represented as follows:

$$P_{\rm CR}=\frac{2\kappa_{\rm B}(p)}{C_{\rm A}^{3/2}\sqrt{at}}\left[\zeta^{5/3}+(D_{\rm NL})^{5/6}\right]^{-3/5}e^{-\zeta^2/4D_{\rm ISM}},\qquad(5)$$

where a is the size of the initial CR cloud, $\zeta=z/\sqrt{C_A at}$, $D_{\rm NL}=F(\Pi)\cdot D_{\rm ISM}\exp(-\Pi)$, with Π being a normalized integrated pressure

$$\Pi=\frac{C_A}{\kappa_B}\int_0^\infty P_{CR}dz.$$

The function F behaves as follows: $F(\Pi)\simeq 2e\approx 5.4$, for $\Pi\gg 1$ and $F(\Pi)\simeq 2\pi\Pi^{-2}$, for $\Pi\ll 1$. Here, $D_{\rm ISM}$ is a normalized background diffusivity $D_{\rm ISM}=\kappa_B/aC_A I_{\rm ISM}$.

To summarize these results, the self-regulated CR escape from a source is characterized by their distribution (partial pressure) comprising the following three zones: (i) a quasi-plateau (core) at small $z/\sqrt{t}<\sqrt{D_{\rm NL}}$ of the height $\sim(D_{\rm NL}t)^{-1/2}$. It is elevated by a factor $\sim\Pi^{-1}\exp(\Pi/2)\gg 1$, compared to the test particle solution because of the strong quasi-linear suppression of the CR diffusion coefficient with respect to its background (test particle) value $D_{\rm ISM}:D_{\rm NL}\sim D_{\rm ISM}\exp(-\Pi)$, (ii) next to the core, where $\sqrt{D_{\rm NL}}<z/\sqrt{t}<\sqrt{D_{\rm ISM}}$, the profile is scale invariant, $P_{\rm CR}\propto 1/z$. The CR distribution in this "pedestal" region is fully self-regulated and independent of Π and $D_{\rm ISM}$ for $\Pi\gg 1$, (iii) the tail of the distribution at $z/\sqrt{t}>\sqrt{D_{\rm ISM}}$ is similar *in shape* to the test particle solution in 1D but it saturates with $\Pi\gg 1$, so that the CR partial pressure is $\propto(D_{\rm ISM}t)^{-1/2}\exp(-z^2/4D_{\rm ISM}t)$, *independent of the strength of the CR source* Π, in contrast to the test-particle regime in which it scales as $\propto\Pi$ ($\Pi\lesssim 1$). Because of the CR diffusivity reduction, the half-life of the CR cloud is increased and its width is decreased, compared to the test particle solution. Depending on the functions $\Pi(p)$ and $D_{\rm ISM}(p)$, the resulting CR spectrum generally develops a spectral break for the fixed values of z and t at the CR momentum p determined by the following relation: $z^2/t\sim D_{\rm NL}(p)\sim D_{\rm ISM}\exp(-\Pi)$.

III. DIFFUSIVE AND HYPERDIFFUSIVE CR TRANSPORT WITH MAGNETIC FOCUSING

Propagating away from their sources, CRs are pitch-angle scattered on weak ISM magnetic irregularities. A seemingly straightforward reduction of kinetic CR description to their spatial transport leads to a diffusive approximation which has the well-known defect of causality violation. There have been attempts at an alternative approach based on the "telegrapher" equation. However, its derivations often lack rigor and transparency and had not been performed to the required (as we show below, fourth order) accuracy. The problem can be formulated very plainly: How to describe CR transport by only their isotropic component, when the anisotropic one is suppressed by the frequent scattering?

The angular distribution of CRs is described by the function $f(\mu,t,z)$[40,41]

$$\frac{\partial f}{\partial t}-\frac{\partial}{\partial\mu}D(\mu)\left(1-\mu^2\right)\frac{\partial f}{\partial\mu}=-\varepsilon\left(\mu\frac{\partial f}{\partial z}+\frac{\sigma}{2}\left(1-\mu^2\right)\frac{\partial f}{\partial\mu}\right).\qquad(6)$$

Here, $\varepsilon=v/l\nu$ is the small parameter of the problem, with v being the particle velocity, l—characteristic scale, and ν—scattering frequency. The dimensionless magnetic mirror inverse scale $\sigma=-B^{-1}\partial B/\partial z$, z points in the local field direction and is measured in the units of l, time in ν^{-1}, and μ is the cosine of the pitch angle, while $D(\mu)\sim 1$ depends on the spectrum of magnetic fluctuations. The isotropic reduction scheme requires a multi-time asymptotic (Chapman-Enskog) expansion. So, we introduce a set of formally independent time variables $t\to t_0,t_1,\ldots$, so that

$$\frac{\partial}{\partial t}=\frac{\partial}{\partial t_0}+\varepsilon\frac{\partial}{\partial t_1}+\varepsilon^2\frac{\partial}{\partial t_2},\ldots,\qquad(7)$$

which leads to the following hierarchy of equations:

$$\begin{aligned}\frac{\partial f_n}{\partial t_0}-\frac{\partial}{\partial\mu}D(\mu)\left(1-\mu^2\right)\frac{\partial f_n}{\partial\mu}&=-\mu\frac{\partial f_{n-1}}{\partial z}-\frac{\sigma}{2}\left(1-\mu^2\right)\\&\quad\times\frac{\partial f_{n-1}}{\partial\mu}-\sum_{k=1}^{n}\frac{\partial f_{n-k}}{\partial t_k}\\&\equiv\mathscr{L}_{n-1}[f](t_0,\ldots,t_n;\mu,z),\end{aligned}\qquad(8)$$

where $f=f_0+\varepsilon f_1+\varepsilon^2 f_2+\cdots$ and the conditions $f_{n<0}=0$ are implied. Using the above expansion, one may obtain an equation for the isotropic part $f_0=\langle f\rangle\equiv(1/2)\int fd\mu$ to arbitrary order in ε. By construction, in no order of approximation will higher time derivatives emerge, as was obviously devised in the Chapman-Enskog method. We terminate this process at the fourth order, ε^4. This is the lowest approximation required to clarify the origin of the telegrapher equation. Higher order terms can, in principle, be calculated at the expense of involved algebra but such calculations would be of no avail. So, our result is as follows:[42]

$$\begin{aligned}\frac{\partial f_0}{\partial t}=\frac{\varepsilon^2}{4}\partial_z'\Big\{&\kappa-\varepsilon\partial_z''\langle\mu W^2\rangle\\&-\frac{\varepsilon^2}{2}\Big[\left(\partial_z''\right)^2\langle W^2(\kappa-U')\rangle\\&-\frac{1}{2}\partial_z'\partial_z\left\langle\frac{\left[\kappa(1-\mu)+U\right]^2}{D(1-\mu^2)}\right\rangle\Big]\Big\}\frac{\partial f_0}{\partial z},\end{aligned}\qquad(9)$$

where $\partial_z' = \partial_z + \sigma$ and $\partial_z'' = \partial_z + \sigma/2$. The coefficients are defined as $\partial W/\partial\mu = 1/D, \langle W\rangle = 0,\ U'(\mu) \equiv \partial U/\partial\mu \equiv (1-\mu^2)/D, U(-1) = 0$, and $\kappa = (1/2)U(1)$.

A. Producing the telegrapher term

Within the employed Chapman-Enskog expansion, the equation for f_0 remains evolutionary in all orders of ε, so no telegrapher term appears. Such term is usually obtained either without clear ordering, e.g., Ref. 43, using specific $D(\mu)$, e.g., Ref. 44, or by truncation of eigenfunction expansion, where the discarded terms may be of the same order in small parameter as those retained, e.g., Refs. 45 and 46. In most of these treatments, care has not been exercised to eliminate the short time scales which are irrelevant to the long-time evolution of the isotropic part of the CR distribution, sought by these reduction schemes. Instead, they retain the second time derivative which changes the type of the resulting transport equation to hyperbolic. As we show below, the second order time derivative term can be recovered from the Chapman-Enskog expansion.

To resolve the above controversy, we simplify Eq. (9) by removing terms unimportant for the controversy. First, we may set $\sigma = 0$ and assume the scattering symmetry, $D(-\mu) = D(\mu)$, to remove the term $\sim\varepsilon^3$, as the $\partial^3/\partial z^3$ term is not included in the telegrapher equation derived for magnetic focusing by, e.g., Ref. 43. Using these simplifications and the slow time $T = \varepsilon^2 t/4$, Eq. (9) rewrites as

$$\frac{\partial f_0}{\partial T} = \kappa\frac{\partial^2 f_0}{\partial z^2} - \varepsilon^2 K\frac{\partial^4 f_0}{\partial z^4}, \tag{10}$$

where K is the hyper-diffusion coefficient

$$K = \frac{1}{2}\left\langle W^2(\kappa - U') - \frac{1}{2}\frac{[\kappa(1-\mu)+U]^2}{D(1-\mu^2)}\right\rangle. \tag{11}$$

To the same order in $\varepsilon \ll 1$, the last equation can be rewritten as follows:

$$\frac{\partial f_0}{\partial T} = \kappa\frac{\partial^2 f_0}{\partial z^2} - \tau\frac{\partial^2 f_0}{\partial T^2}, \tag{12}$$

where $\tau = \varepsilon^2 K/\kappa^2$. This equation has, indeed, the form of a telegrapher equation. However, the comparison of Eq. (10) with, e.g., the telegrapher equation (15) in Ref. 43 shows that the coefficient τ in Eq. (12) is substantially different. The reason is that the equation of Ref. 43 has been obtained by a formal iteration not accounting for all the fourth order terms, the telegrapher term actually originates from. Note that[44] give an expression for τ which is consistent with the result above.[42] More importantly, the telegrapher term in Eq. (12) has a small parameter (at highest derivative). The role of such terms is known from the boundary layer problems. They become crucial near and inside the boundary layer, thus determining its structure and scale. In the context of the telegrapher equation, the boundary layer translates into the initial relaxation phase of the CR distribution. This relaxation is associated with the small scale CR anisotropy in f_n which quickly decays. It should be noted that if a simplified collision term (BGK, or τ-approximation) is used instead of the pitch-angle diffusion in Eq. (6), the telegrapher equation can be accurately derived with no recourse to hyperdiffusion.[47]

To conclude this section, by comparison with the telegrapher equation, the classic Chapman-Enskog is a considerably more suitable and flexible tool to describe the long-time CR propagation, although the telegrapher version (with corrected transport coefficient) may still be useful for studying the magnetically focused CR transport, e.g., Ref. 48. Efforts on improving the CR diffusion models, where their drawbacks are important, need to address the lower level transport, including anisotropic component of the CR distribution, directly. Recent work can be found in, e.g., Ref. 49 and in Sec. IV. Splitting the particle distribution in scattered and unscattered categories is another useful approach, e.g., Refs. 50 and 51. However, when the diffusive treatment is well within the method's validity range (weakly anisotropic spatially smooth CR distributions) neither the telegrapher term nor hyperdiffusivity is essential to the CR transport (see Ref. 42 for more discussion).

IV. SMALL-SCALE CR ANISOTROPY

CR acceleration (e.g., DSA) and propagation models, as discussed in Secs. II and III, typically predict only a large scale, dipolar anisotropy. It would emerge as a small $f_1 \propto \mu$ correction to $f_0 \gg f_1$, produced by localized sources, and can be easily obtained within the treatment outlined in Sec. III. The same is true for the CR self-confinement problem considered in Sec. II, if the small anisotropic correction is taken into account. Now that we expect the CR propagation in the essentially stochastic magnetic fields to be largely ergodic, there is no obvious reason for a significantly sharper than the dipolar anisotropy. Yet observations show that narrow ($\sim 10^\circ$) CR beams do exist.[3–5] They shed new lights on the CR propagation from, and even their acceleration in, putative sources and need to be understood.

A number of scenarios have been suggested to explain the tightly collimated beams. They include magnetic nozzle focusing,[52] propagation effects from local SNR,[53] acceleration in the heliotail,[54,55] and heliosheath propagation effects.[56] Although being plausible, in principle, those explanations impose significant constraints on the relevant parameters and processes. For example, the magnetic mirror ratio must be rather strong to produce $\sim 10^\circ$ anisotropy, and quite a strong magnetic field in the heliotail is required to confine 10 TeV protons and makes the proposed acceleration mechanism work efficiently. Conceptually, different scenarios[57–59] essentially attempt at generating small-scale anisotropy out of the large-scale one by exploiting aspects of interactions between the CRs and MHD turbulence in the ISM. At the first glance, precisely the opposite should occur and the task is clearly of a kind of "squeezing blood out of stone." From a purely mathematical perspective, using certain properties of the particle propagator, these models produce multipoles out of the dipolar component over a long distance (up to a few 100 pc) of particle propagation. At this point, however, the approaches deviate strongly from one another.

In Ref. 59, an interesting technique is employed to generate higher multipoles from the available dipole by using the Liouville's theorem. It is not clear, however, whether the introduction of a simplified collision term in a BGK-form is justified for the treatment of the small-scale anisotropy. The preferred collision operator is the differential one which is much more efficient at smoothing small-scale anisotropies (see, e.g., Sec. III). Attacking the same problem from a different angle, the authors of Ref. 58 rightly state that, although the scattering fields are random, we do not really need to perform an ensemble average, as the current MHD turbulence is static for the limited time observations and fast CRs. There are at least two tests to propose for this explanation. First, as this is actually a magnetic lensing effect with a very long particle path ($L \gg r_g$), small variations of magnetic configuration may produce significant changes in arrival directions of narrow beams. Indeed, the relevant scale of the turbulent field is r_g, so the time scale is $\tau \sim r_g/(V_A + U_{HS})$ with the Alfven and the heliosphere velocities in denominator. The median Milagro energy is $\sim$1 TeV, so for $V_A + U_{HS} \simeq 50$ km/s and $B = 4\mu G$ one obtains $\tau \lesssim 10$ yrs. This may be close enough to the time difference between Milagro and ARGO/HAWC more recent observations. And yes, changes are being observed but they are not quite significant and HAWC is not fully operational yet, so more observations are required and they are underway.[5] The second test has, in fact, already been performed by the authors of Ref. 58. Since CRs interact with the static magnetic fields, their dynamics may be regarded as almost ergodic (strong orbit mixing) on every isoenergetic surface in phase space. Small deviations from ergodicity are responsible for the hot spots in arrival directions. Moving from one energy surface to the next by $\Delta E \sim E$ should strongly decorrelate the spots, since $\Delta r_g \sim r_g$ for them. This is, indeed, observed in simulations carried out in Ref. 58. The upcoming improvements in the energy spectra measurements[5] should substantiate such tests quantitatively and help to discriminate between different mechanisms.

The approach of Ref. 57 is also based on the CR interaction with the ISM turbulence, but includes ensemble averaging, thus removing the above concerns with the short time variability (except for the possible heliospheric variations[60]). The beam direction is assumed to be along the local large scale magnetic field ($l_{loc} \gg r_g$), to minimize the curvature and gradient drifts, that would otherwise evacuate particles from the magnetic tube connecting observer with the source, since the drifts increase with the pitch angle. The following assumptions are made to obtain the beam collimation: (i) *large scale* anisotropic distribution of CRs (generated, for example, by a nearby accelerator, such as a SNR, magnetically connected with the Earth) and (ii) Goldreich-Shridhar[61] (GS) cascade of Alfvenic turbulence originating from a specific scale l, which is the longest scale relevant to the wave-particle interactions.

It is found that the CR distribution develops a characteristic angular shape consisting of a large scale anisotropic part (first eigenfunction of the pitch-angle scattering operator) superposed by a beam, sharply focused in the momentum space along the local field. The large scale anisotropy carries the *momentum dependence* of the source. The following four quantities are tightly constrained by the turbulence scale l: (1) the beam angular width that increases with momentum as $\propto \sqrt{p}$, (2) its fractional excess (with respect to the large scale anisotropic component) that increases as $\propto p$, (3) the maximum momentum, beyond which the beam is destroyed via instability, $p_{\max}$. If the large scale anisotropy originates from a nearby source, magnetically connected with the Earth, the model predicts (4) the range of possible distances to this source, $l_S \sim 100 - 200$ pc. If such source is absent, this range corresponds to the beam collimation length, also a few 100 pc, with the large scale anisotropy originating from the smooth omnigalactic CR gradient. This scale is consistent with the beam collimation length, obtained numerically in Ref. 58.

If the turbulence outer scale l is considered unknown, it can be inferred from any of the first three quantities (1–3) as measured by MILAGRO. All the three quantities consistently imply the same scale $l \simeq 1$ pc. The calculated beam maximum momentum encouragingly agrees with that measured by MILAGRO ($p_{\max} \sim 10$ TeV/c). The theoretical value for the angular width of the beam is found to be $\Delta\vartheta \simeq 4\sqrt{\epsilon}$, where $\epsilon = r_g(p)/l \ll 1$. The beam fractional excess related to the large scale anisotropic part of the CR distribution is $\simeq 50\epsilon$. Both quantities also match the Milagro results near its median energy, that is, for $E \sim 1 - 2$ TeV. So, the beam has a momentum scaling that is one power shallower than the CR carrier, it is drawn from. One interesting conjecture from the $l \simeq 1$ pc requirement is that the proton "knee" at $\simeq 3$ PeV and the beam are of the same origin, as these particles may provide the required outer scale for the MHD turbulence, $r_g \sim l$. Another possibility is to employ the spiral-arm 1-pc value for l, as suggested in Ref. 62.

ACKNOWLEDGMENTS

I am indebted to F. Aharonian, P. Diamond, L. Drury, I. Moskalenko, and R. Sagdeev for collaboration on the topics discussed in this paper. Suggestions by the anonymous referee are also greatly appreciated. This work was supported by the NASA ATP-program under Grant No. NNX14AH36G.

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